

# CONSTRAINTS ON BARYOGENESIS FROM THE DECAY OF SUPERSTRING AXIONS

RAGHAVAN RANGARAJAN\*

*Department of Physics  
University of California  
Santa Barbara, CA 93106*

**ABSTRACT:** We calculate the dilution of the baryon-to-photon ratio by the decay of superstring axions. We find that the dilution is of the order of  $10^7$ . We review several models of baryogenesis and show that most of them can not tolerate such a large dilution. In particular, only one current model of electroweak baryogenesis possibly survives. The Affleck-Dine mechanism in SUSY GUTs is very robust and the dilution by axions could contribute to the dilution required in these models. Baryogenesis scenarios involving topological defects and black hole evaporation are also capable of producing a sufficiently large baryon asymmetry.

(\*: raghu@tpau.physics.ucsb.edu)

*Nuclear Physics B 454 (1995) 369*

In a previous paper, hereafter referred to as Paper I [1], we used cosmological constraints on the decay of axions that arise in  $E'_8 \times E_8$  superstring theories to place a lower limit on the scale of gaugino condensation in the hidden sector. In this paper we calculate the dilution of the baryon-to-photon ratio  $\eta$  by the decay of superstring axions and use this to constrain models of baryogenesis.

In  $E_8 \times E'_8$  models compactified on a Calabi-Yau manifold one typically obtains a model independent axion and several model dependent axions with decay constants related to the compactification scale. If  $E'_8$  breaks down to a non-abelian group there are two non-abelian groups today:  $SU(3)_C$  and the non-abelian subgroup of  $E'_8$ . The model independent and the model dependent axion degrees of freedom rearrange themselves to give two physical axions— the QCD axion  $a$  and the  $E'_8$  axion  $a'$  [2].

The mass, lifetime and energy density of the axion  $a'$  are determined by two energy scales—  $\Lambda$  and the decay constant  $F'_a$ .  $F'_a$  is related to the compactification scale and is about  $10^{15}$  GeV [3].  $\Lambda$  is the scale of gaugino condensation in the hidden sector which breaks supersymmetry. In Paper I, we used limits from nucleosynthesis, limits on the distortion of the cosmic microwave and gamma ray backgrounds and closure arguments to show that the  $E'_8$  axion has to decay before 1s, which implies that  $\Lambda$  has to be greater than  $1.2 \times 10^{13}$  GeV.  $\Lambda$  also determines the masses of the supersymmetric scalars in the observed sector ( $\tilde{m} \sim 10^{-1} \Lambda^3 / M_{Pl}^2$ ). If one assumes that the supersymmetric scalars have masses of the order of  $m_W$  then  $\Lambda$  is approximately  $5 \times 10^{13}$  GeV [4]. It is interesting that our earlier result obtained from astrophysical and cosmological constraints agrees with the independent requirement from particle physics.

In this paper, we let  $\Lambda$  equal  $5 \times 10^{13}$  GeV and study the effects of the decay of the hidden sector axion  $a'$  on the dilution of the baryon-to-photon ratio. For this value of  $\Lambda$ , the hidden sector axion has a mass of  $8.7 \times 10^5$  GeV and a lifetime of  $2.5 \times 10^{-6}$  s. The axion field starts oscillating at a temperature of  $4.5 \times 10^{11}$  GeV and dominates the universe at a temperature of  $1.2 \times 10^5$  GeV ( $1.7 \times 10^{-17}$  s). Its decay leads to a dilution of  $\eta$  by a factor of  $9 \times 10^6$ . The best lower bound on  $\eta$  today, obtained from upper bounds on the  $D + {}^3\text{He}$  abundance, gives  $\eta_f \geq 3 \times 10^{-10}$ . Therefore a successful baryogenesis scenario must produce a baryon-to-photon ratio of  $3 \times 10^{-3}$  or higher. Many current models of baryogenesis are incapable of producing such a large ratio. In the last section we discuss various models and the baryon asymmetry that they can attain.

## II

At high temperatures the  $E'_8$  axion  $a'$  is massless. But as the universe cools it obtains a potential and a mass  $m$ . The axion mass increases till  $T \sim \Lambda$  when it attains its final

low-temperature mass  $m_0$ . For temperatures below  $\Lambda$  the  $a'$  mass is given by [2,5]

$$m_0^2 = \frac{\pi}{450} \frac{\Lambda^6}{M_{Pl}^2 F_a'^2}. \quad (1)$$

For  $\Lambda$  of  $5 \times 10^{13}$  GeV and  $F_a'$  of  $10^{15}$  GeV the low temperature axion mass is  $8.7 \times 10^5$  GeV. In Appendix A of Paper I we show that the universe is radiation dominated when the axion attains its final low temperature mass at  $T \sim \Lambda$ .

When the axion potential appears, the axion field need not be at the minimum of its potential. At a temperature  $T_{\text{osc}}$  the field starts to oscillate about the minimum of the potential with a period  $m^{-1}$ . As we show in Appendix B of Paper I, the axion field starts to oscillate only after it attains its low temperature mass, i.e.,  $T_{\text{osc}} < \Lambda$ . The universe is still radiation dominated when the axion field starts oscillating (Appendix C of Paper I). Eqn. (8) of Paper I gives

$$T_{\text{osc}} = \left[ \frac{m_0 M_{Pl}}{5g_{*_{\text{osc}}}^{1/2}} \right]^{1/2} \quad (2)$$

$g_{*_{\text{osc}}} = 106.75$  is the effective number of relativistic degrees of freedom used to calculate energy density. Therefore  $T_{\text{osc}}$  is  $4.5 \times 10^{11}$  GeV. As the universe cools further, it becomes axion dominated at a temperature  $T_{\text{eq}}$ .  $T_{\text{eq}}$  is  $1.2 \times 10^5$  GeV (Appendix A). Finally, at a temperature  $T_{\text{bd}}$  the axion decays.  $a'$  decays primarily to two gluons through the coupling  $a' F \tilde{F}$ . The gluons create jets of mesons which transfer their energy to the radiation through scattering and annihilations.

The lifetime of the axion is estimated to be

$$\tau = 8200 \pi^5 \frac{F_a'^2}{m_0^3} \quad (3)$$

Thus, the lifetime of the axion is  $2.5 \times 10^{-6}$  s. In Appendix A we show that  $t_{\text{eq}}$  is  $1.7 \times 10^{-17}$  s, i.e., the axion decays only after its energy density dominates the universe.

We study the zero momentum mode of the  $E'_8$  axion field which can be treated as a condensate of zero momentum particles. The axions behave as non-relativistic matter and their energy density falls as  $g_s T^3$ .  $g_s$  is the effective number of relativistic degrees of freedom used to calculate the entropy. Rewriting the energy density in terms of parameters at  $T_{\text{osc}}$  and assuming  $A_{\text{osc}} \sim F_a'$  (see Ref. [6] for a criticism of this assumption) we obtain

$$\rho_a = \frac{5^{3/2}}{2} \left( \frac{g_s}{g_{s_{\text{osc}}}} \right) m_0^{1/2} F_a'^2 \frac{T^3}{M_{Pl}^{3/2}}. \quad (4)$$

$g_{s_{\text{osc}}} = g_{*_{\text{osc}}}$ . Further details of the derivation of the energy density and of the cosmology of superstring axions are given in Paper I.

### III

Now we calculate the dilution of the baryon asymmetry of the universe due to the decay of the axions. We use the approximation that all decays occur at  $t = \tau$  and that the resulting radiation heats up the universe at a temperature  $T_{\text{bd}}$  to a temperature  $T_{\text{ad}}$ . Scherrer and Turner [7] have shown that if one takes into consideration the exponential decay of massive particles, the resulting radiation does not heat up the universe; it causes the universe to cool more slowly to the temperature  $T_{\text{ad}}$ . However, as they point out, their calculation of the change in the entropy is not very different from ours using the naive approximation of instantaneous decay.

The baryon-to-photon ratio  $\eta$  is defined as  $n_B/n_\gamma$ , where  $n_B$  is the baryon number density. We shall assume that the baryon number does not change after baryogenesis. (Any baryon number violating interactions after baryogenesis will only decrease  $n_B$  and strengthen our result.) Therefore,

$$n_{B_i}/n_{B_f} = R_f^3/R_i^3, \quad (5)$$

where  $i$  and  $f$  refer to values at the time of baryogenesis and today respectively and  $R$  is the scale factor.

$$n_{\gamma_f}/n_{\gamma_i} = \frac{T_f^3}{T_i^3} = \frac{S_f}{g_{s_f} R_f^3} \frac{g_{s_i} R_i^3}{S_i} \quad (6)$$

where  $S$  is the entropy.  $g_{s_i}$  and  $g_{s_f}$  are 106.75 and 3.9 respectively. The entropy of the universe does not change when a species disappears through annihilations in equilibrium (such as for  $e^\pm$  annihilations). Therefore,  $S_i = S_{\text{bd}}$  and  $S_f = S_{\text{ad}}$ , where ‘bd’ and ‘ad’ refer to ‘before decay’ and ‘after decay’ respectively.

$$\frac{S_{\text{ad}}}{S_{\text{bd}}} = \frac{g_{s_{\text{ad}}} T_{\text{ad}}^3 R_{\text{ad}}^3}{g_{s_{\text{bd}}} T_{\text{bd}}^3 R_{\text{bd}}^3} \quad (7)$$

In the instantaneous decay approximation  $R_{\text{bd}} = R_{\text{ad}}$  and since the axion is non-relativistic when it decays  $g_{s_{\text{bd}}} = g_{s_{\text{ad}}} = 70$ . Thus,

$$\frac{\eta_i}{\eta_f} = \frac{g_{s_i} T_{\text{ad}}^3}{g_{s_f} T_{\text{bd}}^3} \quad (8)$$

From conservation of energy we get

$$T_{\text{ad}}^3 = \left[ T_{\text{bd}}^4 + \rho_{a_{\text{bd}}} \frac{30}{\pi^2 g_{*_{\text{ad}}}} \right]^{3/4} \quad (9)$$

To obtain  $T_{\text{bd}}$  we first derive the time-temperature relation for  $t_{\text{eq}} < t < \tau$ . For an  $\Omega = 1$  universe,

$$H = \left( \frac{8\pi}{3M_{Pl}^2} \rho \right)^{1/2} \quad (10)$$

When the universe is axion dominated we use the axion energy density  $\rho_a$  in (4). Also, for  $t \gg t_{\text{eq}}$ ,  $H = (2/3t)$ . Setting  $t = \tau$  we get

$$T_{\text{bd}} = 0.21 \left[ \frac{g_{s_{\text{osc}}}^{1/4} M_{Pl}^{7/2}}{g_{s_{\text{bd}}} \tau^2 m_0^{1/2} F_a'^2} \right]^{1/3} \quad (11)$$

We also get  $\rho_{a_{\text{bd}}}$  from (10).

$$\rho_{a_{\text{bd}}} = \frac{M_{Pl}^2}{6\pi\tau^2} \quad (12)$$

Thus  $T_{\text{bd}}$  is 5.8 MeV and  $T_{\text{ad}}$  is 0.39 GeV.

Substituting (9), (11) and (12) in (8) we get

$$\frac{\eta_i}{\eta_f} = \frac{g_{s_i}}{g_{s_f}} \left[ 1 + \left( \frac{M_{Pl}^2}{6\pi\tau^2} \right) \left( \frac{30}{\pi^2 g_{*_{\text{ad}}}} \right) \left( 110 \frac{g_{s_{\text{bd}}}^{1/4} \tau^2 m_0^{1/2} F_a'^2}{g_{s_{\text{osc}}}^{1/4} M_{Pl}^{7/2}} \right)^{(4/3)} \right]^{3/4} \quad (13)$$

Since the second term in the bracket is much larger than 1, or equivalently,  $\rho_{a_{\text{bd}}} \gg \rho_{\text{rad}}$  at the time of the decay

$$\frac{\eta_i}{\eta_f} = 28 \frac{g_{s_i}}{g_{s_f}} \frac{g_{s_{\text{bd}}}}{g_{s_{\text{osc}}}^{1/4} g_{*_{\text{ad}}}^{3/4}} \frac{\tau^{1/2} m_0^{1/2} F_a'^2}{M_{Pl}^2} \quad (14)$$

Substituting for  $\tau$  and  $m_0$  in terms of  $\Lambda$  we get

$$\frac{\eta_i}{\eta_f} = 5.3 \times 10^5 \frac{g_{s_i}}{g_{s_f}} \frac{g_{s_{\text{bd}}}}{g_{s_{\text{osc}}}^{1/4} g_{*_{\text{ad}}}^{3/4}} \frac{F_a'^4}{\Lambda^3 M_{Pl}} \quad (15)$$

For  $\Lambda$  of  $5 \times 10^{13}$  GeV and  $F_a'$  of  $10^{15}$  GeV, the dilution is

$$\frac{\eta_i}{\eta_f} = 9 \times 10^6 \quad (16)$$

For a conservative lower bound on  $\eta_f$  of  $3 \times 10^{-10}$ , we find that  $\eta_i$  must be greater than  $3 \times 10^{-3}$  or  $n_B/s$ , the baryon-to-entropy ratio, must be greater than  $10^{-5}$ .  $n_B/s$  today is  $4 \times 10^{-11}$ .

## IV

### Discussion:

The simplest mechanism for producing a baryon asymmetry involves out-of-equilibrium decays of massive Higgs or gauge bosons in GUTs or supersymmetric GUTs. This mechanism can, in principle, produce a baryon-to-entropy ratio of  $10^{-4}$  [8]. However, the maximum asymmetry obtained in most specific GUT scenarios is the observed asymmetry today or less [9]. In general, one obtains an even smaller asymmetry in supersymmetric GUTs. In the context of an inflationary universe, out-of-equilibrium decays in GUTs fare worse [10]. The COBE results constrain the inflationary potential and prescribes a low inflaton mass  $\sim 10^{11}$  GeV. The out-of-equilibrium decays scenarios then require the massive boson to be lighter than the inflaton which makes it difficult to obtain even the present asymmetry. Again, the situation is much worse for SUSY GUTs.

However, the Affleck-Dine mechanism involving the decay of sfermion condensates in SUSY GUTs in an inflationary universe can produce a baryon-to-entropy ratio as large as  $10^{-2}$  [11]. Furthermore, Davidson et al. [12] have shown that the presence of Bose-Einstein condensates can suppress the destruction of the baryon asymmetry by electroweak sphaleron processes. To preserve a Bose-Einstein condensate till the electroweak phase transition in an inflationary universe, below which sphaleron processes are naturally suppressed, requires  $n_B/s \geq 0.01$ . It will subsequently need to be diluted to about  $4 \times 10^{-11}$  before nucleosynthesis. Superstring axion decays can certainly provide part of the required dilution. (For other mechanisms to dilute the entropy see ref. [13].)

Fukugita and Yanagida [14] (also see ref. [15]), Lazarides and Shafi [16] and Campbell et al. [17,10] produce a baryon asymmetry by first creating a lepton asymmetry and then convert this into a baryon asymmetry by sphaleron processes. In the Fukugita-Yanagida and Lazarides-Shafi mechanism, a lepton asymmetry is created by the out-of-equilibrium decay of heavy right-handed Majorana neutrinos in a see-saw mechanism. In the model of Lazarides and Shafi the heavy neutrinos are obtained by the decay of the inflaton field. Campbell et al. [10] consider a supersymmetric version of the Fukugita-Yanagida model and also study it within the context of an inflationary universe. Though these models can produce the asymmetry observed today they can not create an asymmetry as high as  $10^{-5}$ . Campbell et al. also create a lepton asymmetry by the effect of lepton number violating induced operators, arising from see-saw (s)neutrino masses, which act on scalar condensate oscillations along flat directions of the supersymmetric standard model. As in the Affleck-Dine mechanism, a large asymmetry of  $10^{-5}$  can be created.

In recent years, there have been a number of models attempting to create the baryon asymmetry at the electroweak transition, generally in a first-order phase transition. In a

model by Nelson et al. [18], a hypercharge asymmetry is created outside a bubble of true vacuum by a difference in the reflection rates for left- and right-handed top quarks bouncing off the bubble wall in the false vacuum. The hypercharge asymmetry is converted into a baryon asymmetry. They can obtain a baryon-to-entropy ratio of  $10^{-6}$ . However that assumes maximal CP violation which is not very likely [19]. In a similar scenario, a lepton asymmetry is created by the reflection of heavy neutrinos off the bubble wall [20]. The lepton asymmetry is then converted into a baryon asymmetry by B+L violating processes. This scenario can produce the required  $n_B/s$  of  $10^{-5}$  if there is a fourth neutrino heavier than 45 GeV (as required by the data from Z decays) and the neutrino mixing angles are small. However this value may be diminished by diffusion of particles from the true vacuum to the false vacuum [21]. Furthermore, one has to check that the heavy neutrino spends enough time in the symmetric phase after reflection for sphaleron processes to convert the lepton asymmetry into a baryon asymmetry, before the wall catches up with it [19].

Turok and Zadrozny [22] and McLerran et al. [23] produce a baryon asymmetry through the CP asymmetric interaction of the wall with field fluctuations of non-zero Chern-Simons number in the the symmetric phase [24]. Sufficient CP violation is typically provided by extending the Standard Model to two or more Higgs doublets. The above scenarios can produce a baryon asymmetry that is close to the present value only. (Dine et al. [25] have indicated that the baryon asymmetry produced will be even lower.)

Another mechanism for producing the baryon asymmetry at the electroweak phase transition has been proposed by McLerran [26]. In this model CP violation is provided by the existence of a QCD axion. Interference between electroweak sphaleron-induced baryon number violating processes and QCD sphaleron-induced CP violating processes (which require  $\theta_{\text{QCD}} \neq 0$ ) produces a baryon asymmetry. In the standard model this scenario can not produce the present day asymmetry. Introducing heavy scalars in the theory can give a maximum  $n_B/s$  of  $10^{-8}$ .

Theories of spontaneous baryogenesis involve the spontaneous breaking of a  $U(1)_B$  baryon symmetry [27]. Regions of baryons and anti-baryons, inflated to supra-horizon sizes, are produced by the slow rolling of the consequent pseudo-Goldstone boson towards the minimum of its potential and by its decay as it finally oscillates about its minimum. While such scenarios can produce the observed asymmetry they can not tolerate a dilution of the order of  $10^5$ . This idea of creating a baryon asymmetry by using the evolution of a scalar to create an effective chemical potential has been extended to the electroweak phase transition by Dine et al. [28,25] and by Cohen et al. [29] and Dine and Thomas [30]. The maximum baryon-to-entropy ratio they obtain is of the order of  $10^{-7}$ .

Topological defects have also been used to create a baryon asymmetry. Nussinov [31]

has considered monopole-anti-monopole annihilations as a source of heavy GUT Higgs and gauge bosons which decay asymmetrically into baryons and anti-baryons giving a baryon asymmetry. This model can produce a maximum baryon-to-entropy ratio of  $10^{-10}$ . Bhattacherjee et al. [32] can obtain a baryon asymmetry of  $10^{-5}$  from the decay of heavy Higgs and gauge bosons produced by collapsing cosmic string loops if the GUT scale is about  $10^{16}$  GeV. Brandenberger et al. [33] have also obtained a high baryon asymmetry similarly from collapsing string loops. Kawasaki and Maeda [34] consider a baryon asymmetry created by the decay of particles emitted from cusps moving close to the speed of light on cosmic strings, as well as from kinky string loops, proposed by Garfinkle and Vachaspati [35]. They can create a large enough asymmetry but the existence of cusp evaporation is uncertain. Mohazzab [36] also obtains a baryon asymmetry as high as  $10^{-5}$  from cusp annihilation on long cosmic strings. Barrow et al. [37] create a baryon asymmetry from the decay of bosons produced in bubble wall collisions in extended inflation models. They can obtain a baryon asymmetry of  $10^{-5}$  only if the colliding bubbles are not much larger than the critical radius.

Several models have been proposed in which a baryon asymmetry is created by the evaporations of black holes. Zel'dovich [38] and Dolgov [39] showed that one can obtain a baryon asymmetry from black holes even if baryonic charge is conserved. Though initially the black hole emits particles that decay to an equal number of baryons and antibaryons, more antibaryons than baryons are recaptured by the black hole. Toussaint et al. [40], Turner and Schramm [41], Turner [42], Barrow [43] and Barrow and Ross [44] have suggested that black holes can emit heavy bosons and produce a baryon asymmetry by their decay. These models require baryon number violation and are similar to the GUT baryogenesis scenarios. Barrow et al. [45] have obtained a baryon asymmetry from the evaporation of black holes formed during extended inflation. Black hole evaporation can give a very large baryon asymmetry. However it is difficult to estimate their number density and mass distribution. Hence there is large uncertainty in the baryon asymmetry that can actually be created.

**Conclusion:** Thus we see that the existence and decay of superstring axions is not compatible with many models of baryogenesis. Only one current model of electroweak baryogenesis may be able to tolerate the dilution of the baryon asymmetry. The Affleck-Dine mechanism is certainly very robust. Models involving topological defects and black hole evaporation are also capable of creating a large enough baryon asymmetry.

**Acknowledgements:** I would like to thank Mark Srednicki, Subir Sarkar, Robert Scherrer, David Kaplan and Kiwoon Choi for very useful discussions. I would also like to thank the referee for pointing out the possibility of the induced decay of axions due to res-



onance effects. I am also grateful to the Center for Particle Astrophysics at the University of California, Berkeley, where part of this work was completed, for their hospitality.

After this work was completed, I discovered that similar work on the cosmological consequences of scalars in superstring and supergravity theories has been carried out by other authors [46].

This work was supported by NSF Grant No. PHY91-16964.

## Appendix A

In this Appendix we show that the axion decays after its energy density dominates the universe.

At  $t = t_{\text{eq}}$ ,

$$(4t_{\text{eq}}^2)^{-1} = \frac{8\pi}{3M_{Pl}^2} \left( \frac{1}{2} m_0^2 A^2(T_{\text{eq}}) \right) \quad (A.1)$$

Also,

$$m_0 A^2(T_{\text{eq}}) = m_0 A^2(T_{\text{osc}}) \frac{T_{\text{eq}}^3}{T_{\text{osc}}^3} \quad (A.2)$$

where we have set  $g_{s_{\text{osc}}} = g_{s_{\text{eq}}} = 106.75$ .  $T_{\text{osc}}$  is given by (2). We now calculate  $T_{\text{eq}}$ .

At  $T_{\text{eq}}$ ,

$$\begin{aligned} \frac{\pi^2}{30} g_{*_{\text{eq}}} T_{\text{eq}}^4 &= \left( \frac{1}{2} m_0^2 A^2(T_{\text{eq}}) \right) \\ &= \left( \frac{\pi}{900} \frac{\Lambda^6}{M_{Pl}^2 F_a'^2} \right) A^2(T_{\text{osc}}) \frac{T_{\text{eq}}^3}{T_{\text{osc}}^3} \end{aligned} \quad (A.3)$$

Therefore,

$$T_{\text{eq}} = \frac{1.1 \times 10^{-2}}{g_{*_{\text{eq}}}} \frac{\Lambda^6}{M_{Pl}^2 T_{\text{osc}}^3} \quad (A.4)$$

$$= 1.5 \frac{\Lambda^{3/2} F_a'^{3/2}}{M_{Pl}^2} \quad (A.5)$$

$$T_{\text{eq}} = 1.2 \times 10^5 \text{ GeV}.$$

Combining (A.1), (A.2), (A.5) and (2), we get

$$\begin{aligned} t_{\text{eq}} &= 1.3 \times 10^{-2} \frac{M_{Pl}^5}{\Lambda^3 F_a'^3} \\ &= 1.7 \times 10^{-17} \text{ s} \end{aligned} \quad (A.6)$$

However  $\tau = 8200\pi^5 (F_a'^2/m_0^3)$  gives a lifetime of  $2.5 \times 10^{-6}$  s. Thus the axion decays after it has dominated the universe, or  $T_{\text{eq}} > T_{\text{bd}}$ .

## REFERENCES

- [1] R. Rangarajan, Nucl. Phys. B454 (1995) 357.

- [2] K. Choi and J. E. Kim, Phys. Lett. B165 (1985) 71.
- [3] J. E. Kim, Phys. Rep. 150 (1987) 1.
- [4] J. P. Derendinger, L. E. Ibanez and H. P. Nilles, Phys. Lett. B155 (1985) 65.
- [5] M. Dine, R. Rohm, N. Seiberg and E. Witten, Phys. Lett. B156 (1985) 55.
- [6] A. D. Linde, Phys. Lett. B201 (1988) 437; Phys. Lett. B259 (1991) 38.
- [7] R. J. Scherrer and M. S. Turner, Phys. Rev. D 31 (1985) 681.
- [8] E. W. Kolb and M. S. Turner, Ann. Rev. Nucl. Part. Sci. 33 (1983) 645.
- [9] A. D. Dolgov, Phys. Rep. 222 (1992) 309. This is an excellent review of baryogenesis models.
- [10] B. Campbell, S. Davidson and K. A. Olive, Nucl. Phys. B399 (1993) 111.
- [11] I. Affleck and M. Dine, Nucl. Phys. B249, (1985) 361; A. D. Linde, Phys. Lett. B160, (1985) 243.
- [12] S. Davidson, H. Murayama and K. A. Olive, hep-ph 9403259, (1994).
- [13] K. Yamamoto, Phys. Lett. B168, (1986) 341; K. Enqvist, D. V. Nanopoulos and M. Quiros, Phys. Lett. B169, (1986) 343; J. Ellis, K. Enqvist, D. V. Nanopoulos and K. A. Olive; Phys. Lett. B188, (1987) 415.
- [14] M. Fukugita and T. Yanagida, Phys. Lett. B174 (1986) 45; Phys. Rev. D 42 (1990) 1285.
- [15] M. Luty, Phys. Rev. D 45 (1992) 455.
- [16] G. Lazarides and Q. Shafi, Phys. Lett. B258 (1991) 305.
- [17] B. Campbell, S. Davidson and K. A. Olive, Phys. Lett. B303 (1993) 63.
- [18] A. E. Nelson, D. B. Kaplan and A. G. Cohen, Nucl. Phys. B373 (1992) 453.
- [19] D. B. Kaplan, private communication.
- [20] A. G. Cohen, D. B. Kaplan and A. E. Nelson, Phys. Lett. B245 (1990) 561; Nucl. Phys. B349 (1991) 727.
- [21] J. Maalampi, J. Sirrka and I. Vilja, hep-ph 9405378, (1994).
- [22] N. Turok and J. Zadrozny, Phys. Rev. Lett. 65 (1990) 2331; Nucl. Phys. B358 (1991) 471.
- [23] L. McLerran, M. Shaposhnikov, N. Turok and M. Voloshin, Phys. Lett. B256 (1991) 451.
- [24] J. Ambjorn, M. Laursen and M. E. Shaposhnikov, Phys. Lett. B197 (1987) 49; Nucl. Phys. B316 (1989) 483.

- [25] M. Dine, P. Huet and R. Singleton Jr., Nucl. Phys. B375 (1992) 625.
- [26] L. McLerran, Phys. Rev. Lett. 62 (1989) 1075.
- [27] A. G. Cohen and D. B. Kaplan, Phys. Lett. B199 (1987) 251; Nucl. Phys. B308 (1988) 913.
- [28] M. Dine, P. Huet, R. Singleton Jr. and L. Susskind, Phys. Lett. B257 (1991) 351.
- [29] A. G. Cohen, D. B. Kaplan and A. E. Nelson, Phys. Lett. B263 (1991) 86; hep-ph 9406345, (1994).
- [30] M. Dine and S. Thomas, hep-ph 9401265, (1994).
- [31] S. Nussinov, Phys. Lett. B110 (1982) 221.
- [32] R. Bhattacharjee, T. W. Kibble and N. Turok, Phys. Lett. B119 (1982) 95.
- [33] R. H. Brandenberger, A. C. Davis and M. Hindmarsh, Phys. Lett. B263 (1991) 239.
- [34] M. Kawasaki and K. Maeda, Phys. Lett. B208 (1988) 84.
- [35] D. Garfinkle and T. Vachaspati, Phys. Rev. D 36 (1987) 2229.
- [36] M. Mohazzab, hep-ph 9409274, (1994).
- [37] J. D. Barrow, E. J. Copeland, E. W. Kolb and A. R. Liddle, Phys. Rev. D 43 (1991) 977.
- [38] Ya. B. Zel'dovich, JETP Lett. 24 (1976) 24.
- [39] A. D. Dolgov, Phys. Rev. D 24 (1981) 1042.
- [40] D. Toussaint, S. B. Treiman, F. Wilczek, A. Zee, Phys. Rev. D 19 (1979) 1036.
- [41] M. S. Turner and D. N. Schramm, Nature 279 (1979) 303.
- [42] M. S. Turner, Phys. Lett. B89 (1979) 155.
- [43] J. D. Barrow, Mon. Not. R. Astr. Soc. 192 (1980) 427.
- [44] J. D. Barrow and G. G. Ross, Nucl. Phys. B181 (1981) 461.
- [45] J. D. Barrow, E. J. Copeland, E. W. Kolb and A. R. Liddle, Phys. Rev. D 43 (1991) 984.
- [46] G. D. Coughlan, W. Fischler, E. W. Kolb, S. Raby and G. G. Ross, Phys. Lett. B131 (1983) 59; A. S. Goncharov, A. D. Linde and M. I. Vysotsky, Phys. Lett. B147 (1984) 279; G. German and G. G. Ross, Phys. Lett. B172 (1986) 305; J. Ellis, D. Nanopoulos and M. Quiros, Phys. Lett. B174 (1986) 176; O. Bertolami, Phys. Lett. B209 (1988) 277.